

Natural Chaotic Inflation in Supergravity and Leptogenesis

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Abstract

We comprehensively investigate a chaotic inflation model proposed recently in the framework of supergravity. In this model, the form of Kähler potential is determined by a symmetry, that is, the Nambu-Goldstone-like shift symmetry, which guarantees the absence of the exponential factor in the potential for the inflaton field. Though we need the introduction of small parameters, the smallness of the parameters is justified also by symmetries. That is, the zero limit of the small parameters recovers symmetries, which is natural in the 't Hooft's sense. The leptogenesis scenario via the inflaton decay in this chaotic inflation model is also discussed. We find that the lepton asymmetry enough to explain the present baryon number density is produced for low reheating temperatures avoiding the overproduction of gravitinos.

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I. INTRODUCTION

Big-bang cosmology is a very attractive theory because it explains well the three main observational results in cosmology, that is, Hubble expansion, the cosmic microwave background radiation (CMBR), and the primordial abundance of light elements. But it has famous problems, namely, the horizon problem and the flatness problem, and does not account for the origin of primordial fluctuations of CMBR as observed by the Cosmic Background Explorer (COBE) satellite [1]. The most natural solution to these problems is inflation [2]. Until now, many types of inflation models have been proposed. Among them, chaotic inflation is special in that it can take place at about the Planck time. Other types of inflation occur generally at much later times so that they suffer from the flatness (longevity) problem [2] though it is milder than the original one, that is, why the universe lives so long up to the low energy scale. Furthermore, other types of inflation except chaotic and topological inflation also suffer from the initial value problem [2,4], that is, why the inflaton field φ is homogeneous over the horizon scale and lies in the small region of the potential which leads to a successful inflation. If the universe is open at the beginning [3], the flatness problem may be evaded and the topological inflation may occur.¹ However, the chaotic inflation gives the most natural solution to the above problems since it takes place at about the Planck time. Thus, the chaotic inflation is the most attractive inflation without any fine tuning.

The fact that inflation takes place at higher energy scales than the electroweak scale confronts us with a hierarchy problem between such two energy scales. One of the most attractive solutions is supersymmetry (SUSY) [5], which stabilizes such a large hierarchy against radiative corrections. Thus, it is important to consider inflation in the framework of the local version of SUSY, i.e., supergravity.

Chaotic inflation can be realized for a very simple polynomial potential. Due to this simplicity, a lot of applications have been investigated for the chaotic inflation, for example, preheating [6], superheavy particle production [7], and primordial gravitational waves [8]. It is, however, very difficult to realize such a polynomial potential in supergravity because the minimal supergravity potential has an exponential factor ($e^{\varphi^\dagger\varphi/M_G^2+\dots}$), which prevents inflaton φ from having an initial value much larger than the gravitational scale $M_G \simeq 2.4 \times 10^{18}$ GeV. Thus, it has been believed to be very difficult in incorporating the chaotic inflation in the framework of supergravity. Although some models for the chaotic inflation were proposed using specific Kähler potentials instead of the canonical Kähler potential [9,10], such Kähler potentials have no symmetry reason and we must invoke a fine tuning.

However, we have recently constructed a natural chaotic inflation model in supergravity without any fine tuning [11]. The term “natural” has two meanings. First of all, the form of Kähler potential is determined by a symmetry, that is, the Nambu-Goldstone-like shift symmetry, which guarantees the absence of the exponential factor in the potential for the inflaton field. Though we need the introduction of small breaking parameters, the smallness of parameters is justified also by symmetries. That is, the zero limit of small parameters

¹Exactly speaking, for a successful topological inflation in supergravity, the Kähler potential must be fine-tuned against quantum corrections in order to keep the flatness of the potential near the origin.

recovers symmetries, which is natural in the 't Hooft's sense [12]. This is the second meaning of our term “natural.” In this paper, we comprehensively investigate this chaotic inflation model, particularly paying attention to the small parameters of symmetry breaking in the superpotential.

As an application of the above new type of chaotic inflation model [11], we discuss the leptogenesis. Recent experimental results on the atmospheric neutrinos strongly indicate that neutrinos have small masses of the order of 0.01–0.1 eV [13]. Such small masses are naturally explained by the seesaw mechanism [14], which predicts superheavy right-handed neutrinos. The presence of Majorana masses of right-handed neutrinos naturally leads to the leptogenesis because it violates the lepton number conservation. The decay of superheavy Majorana neutrinos produces the lepton number asymmetry, in particular, $B-L$ asymmetry if C and CP symmetries are broken, which is converted into baryon asymmetry [15] through the sphaleron effects [16]. Therefore, we discuss a leptogenesis scenario in the above mentioned chaotic inflation model.

In the next section, we briefly review on the chaotic inflation model in supergravity. In Sec. III, we investigate the dynamics of chaotic inflation. In Sec. IV, we discuss the leptogenesis via the inflaton decay. The last section is devoted to discussion and conclusions.

II. NATURAL CHAOTIC INFLATION MODEL IN SUPERGRAVITY

As explained in the introduction, the chaotic inflation is special in that it takes place around the gravitational scale and hence it does not suffer from the flatness (longevity) and the initial value problems. But it was a long-standing problem to realize a chaotic inflation naturally in supergravity because the minimal supergravity potential has an exponential growth ($e^{\varphi^\dagger\varphi/M_G^2+\dots}$) for the inflaton field φ , which prevents the inflaton φ from taking an initial value much larger than the gravitational scale. However, we have recently proposed a natural chaotic inflation model in supergravity by imposing Nambu-Goldstone-like shift symmetry. In this section, we briefly review our chaotic inflation model [11].

For the inflaton chiral supermultiplet $\Phi(x, \theta)$, we assume that the Kähler potential $K(\Phi, \Phi^*)$ is invariant under the shift of Φ ,²

$$\Phi \rightarrow \Phi + i C M_G, \quad (1)$$

where C is a dimensionless real parameter. Hereafter, we set M_G to be unity. Thus, the Kähler potential is a function of $\Phi + \Phi^*$, i.e. $K(\Phi, \Phi^*) = K(\Phi + \Phi^*)$. It is now clear that the supergravity effect $e^{K(\Phi+\Phi^*)}$ discussed above does not prevent the imaginary part of the scalar components of Φ from having a value larger than the gravitational scale. So, we identify it with the inflaton field φ [see Eq. (8)]. As long as the shift symmetry is exact, the inflaton φ never has a potential and hence it never causes inflation. Therefore, we need some breaking term in the superpotential. Here, we discuss the form of the superpotential.

²The inflaton Φ may be one of modulus fields in string theories. We hope that the explicit breaking of the shift symmetry introduced below will be understood by yet unknown dynamics of string theories.

First of all, we assume that in addition to the shift symmetry, the superpotential is invariant under the $U(1)_R$ symmetry, which prohibits a constant term in the superpotential. Then, the above Kähler potential is invariant only if the R charge of Φ is zero. Therefore, the superpotential comprised of only the Φ field is not invariant under the $U(1)_R$ symmetry, which compels us to introduce another supermultiplet $X(x, \theta)$ with its R-charge equal to two.

We now introduce a suprior field Ξ describing the breaking of the shift symmetry, and extend the shift symmetry including the suprior field Ξ as follows,³

$$\begin{aligned}\Phi &\rightarrow \Phi + i C, \\ \Xi &\rightarrow \frac{\Phi}{\Phi + i C} \Xi.\end{aligned}\tag{2}$$

That is, the combination $\Xi\Phi$ is invariant under the shift symmetry. Then, the general superpotential invariant under the shift and $U(1)_R$ symmetries is given by

$$W = X \left\{ \Xi\Phi + \alpha_3(\Xi\Phi)^3 + \dots \right\} + \delta_1 X \left\{ 1 + \alpha_2(\Xi\Phi)^2 + \dots \right\},\tag{3}$$

where we have assumed the R charge of Ξ vanish. The shift symmetry is softly broken by inserting the vacuum value $\langle \Xi \rangle = m$. The mass parameter m is fixed at a value much smaller than unity representing the magnitude of breaking of the shift symmetry (2). We see that higher order terms with α_i of the order of unity become irrelevant for the dynamics of the chaotic inflation. Thus, we neglect them in the following discussion⁴ unless explicitly mentioned. We should note that the complex constant δ_1 is also of the order of unity in general. But, as shown later, the absolute magnitude of δ_1 must be at most of the order of m , which is much smaller than unity. Therefore, we introduce the Z_2 symmetry, under which both the Φ and X fields are odd. Then, the smallness of the constant δ_1 is associated with the small breaking of the Z_2 symmetry. That is, we introduce a suprior field Π with odd charge under the Z_2 symmetry. The vacuum value $\langle \Pi \rangle = \delta_1$ breaks the Z_2 symmetry. Though the above superpotential is not invariant under the shift and the Z_2 symmetries, the model is completely natural in the 't Hooft's sense [12] because we have enhanced symmetries in the limit m and $\delta_1 \rightarrow 0$. We use, in the following analysis, the superpotential,

$$W \simeq mX\Phi + \delta_1 X.\tag{4}$$

The Kähler potential invariant under the shift and $U(1)_R$ symmetries is give by

$$K = \delta_2(\Phi + \Phi^*) + \frac{1}{2} (\Phi + \Phi^*)^2 + XX^* + \dots.\tag{5}$$

Here $\delta_2 \sim |\delta_1|$ is a real constant representing the breaking effect of the Z_2 symmetry. The terms $\delta_3 m_3 \Phi + \delta_3^* m_3^* \Phi^*$ and $(m_4 \Phi)^2 + (m_4^* \Phi^*)^2$ may appear, where δ_3 and m_4 are complex

³If Ξ transforms as $\Xi \rightarrow \frac{\Phi^n}{(\Phi + iC)^n} \Xi$ ($n \geq 2$), we have $W = X\Xi\Phi^n$, which may cause φ^{2n} chaotic inflations.

⁴Among all complex constants, only a constant becomes real by use of the phase rotation of the X field. Below we set m to be real.

constants representing the breaking of the Z_2 and the shift symmetries ($|\delta_3| \sim |\delta_1|$ and $|m_3| \sim |m_4| \sim m$). But, these terms are extremely small so we have omitted them in the Kähler potential (5). We have also omitted a constant term because it only changes the overall factor of the potential, whose effect can be renormalized into the constant m and δ_1 . Here and hereafter, we use the same characters for scalar with those for corresponding supermultiplets.

III. DYNAMICS OF CHAOTIC INFLATION

The Lagrangian density $\mathcal{L}(\Phi, X)$ neglecting the higher order terms is given by

$$\mathcal{L}(\Phi, X) = \partial_\mu \Phi \partial^\mu \Phi^* + \partial_\mu X \partial^\mu X^* - V(\Phi, X), \quad (6)$$

with the potential $V(\Phi, X)$,

$$V(\Phi, X) = m^2 e^K \left[|X|^2 |1 + (\delta_2 + \Phi + \Phi^*)(\Phi + \delta'_1)|^2 + |\Phi + \delta'_1|^2 (1 - |X|^2 + |X|^4) \right], \quad (7)$$

with $\delta'_1 \equiv \delta_1/m$. Now, we decompose the complex scalar field Φ into two real scalar fields as

$$\Phi = \frac{1}{\sqrt{2}}(\eta + i\varphi), \quad (8)$$

where we identify φ with the inflaton. Then, the Lagrangian density $\mathcal{L}(\eta, \varphi, X)$ is given by

$$\mathcal{L}(\eta, \varphi, X) = \frac{1}{2} \partial_\mu \eta \partial^\mu \eta + \frac{1}{2} \partial_\mu \varphi \partial^\mu \varphi + \partial_\mu X \partial^\mu X^* - V(\eta, \varphi, X), \quad (9)$$

with the potential $V(\eta, \varphi, X)$,

$$\begin{aligned} V(\eta, \varphi, X) = & m^2 e^{-\frac{\delta_2^2}{2}} \exp \left\{ \left(\eta + \frac{\delta_2}{\sqrt{2}} \right)^2 + |X|^2 \right\} \\ & \times \left[|X|^2 \left\{ 1 + 2 \left(\eta + \frac{\delta_2}{\sqrt{2}} \right) (\eta + \delta_R) + \left(\eta + \frac{\delta_2}{\sqrt{2}} \right)^2 [(\eta + \delta_R)^2 + (\varphi + \delta_I)^2] \right\} \right. \\ & \left. + \frac{1}{2} \{ (\eta + \delta_R)^2 + (\varphi + \delta_I)^2 \} (1 - |X|^2 + |X|^4) \right]. \end{aligned} \quad (10)$$

Here, the complex constant δ'_1 is decomposed into a real and an imaginary part,

$$\delta'_1 = \frac{1}{\sqrt{2}}(\delta_R + i\delta_I). \quad (11)$$

Note that η and $|X|$ should be taken as $|\eta|, |X| \lesssim \mathcal{O}(1)$ for $\delta_2 \ll 1$ because of the presence of e^K factor. On the other hand, φ can take a value much larger than $\mathcal{O}(1)$ since e^K does not contain φ . For the case $\eta, |X| \ll \mathcal{O}(1)$, which is valid during the inflation as shown later, the potential can be approximated as

$$V(\eta, \varphi, X) \simeq \frac{1}{2} m^2 \tilde{\varphi}^2 + m^2 |X|^2, \quad (12)$$

where $\tilde{\varphi} \equiv \varphi + \delta_I$ and we have taken $me^{-\delta_2^2/4} \simeq m$ since $\delta_2 \ll 1$.

Thus, the term proportional to $\tilde{\varphi}^2$ becomes dominant and the chaotic inflation takes place if the initial value $\tilde{\varphi} \gg 1$. The potential minimum for η during the inflation, η_m , is given by the minimum of the Kähler potential, which yield $\eta_m \simeq -\frac{\delta_2}{\sqrt{2}}$. Then, during the chaotic inflation, the effective mass squared of η , m_η^2 , becomes

$$m_\eta^2 \simeq m^2 \tilde{\varphi}^2 \simeq 6H^2, \quad (13)$$

where $H[\simeq \frac{1}{\sqrt{6}}m\tilde{\varphi}]$ is the Hubble parameter. Because m_η^2 is much larger than $\frac{9}{4}H^2$, the field η rapidly oscillates around the minimum η_m with its amplitude damped in proportion to $a^{-3/2}$, where a is the scale factor. Thus, the field η settles down to the minimum η_m very quickly.

On the other hand, the effective mass of X , m_X , is m , which is smaller than the Hubble scale so that it does not oscillate but only slow rolls.⁵ Using the slow-roll approximation, the classical equations of motion for both $\tilde{\varphi}$ and X fields are given by

$$3H\dot{\tilde{\varphi}} \simeq -m^2\tilde{\varphi}, \quad (14)$$

$$3H\dot{X} \simeq -m^2X, \quad (15)$$

where the overdot represents the time derivative. Also, here and hereafter, we assume that X is real and positive. Then, we obtain the relation between $\tilde{\varphi}$ and X fields,

$$\left(\frac{X}{X(0)}\right) \simeq \left(\frac{\tilde{\varphi}}{\tilde{\varphi}(0)}\right), \quad (16)$$

where $\tilde{\varphi}(0)$ and $X(0)$ are the initial values of $\tilde{\varphi}$ and X fields. But, one should note that this relation actually holds if and only if quantum fluctuations are unimportant for both $\tilde{\varphi}$ and X fields. Therefore, we need to clarify when the classical description is feasible. For this purpose, we first compare quantum fluctuations with classical changes for the field $\tilde{\varphi}$. During one expansion time, by use of Eqs. (13) and (14), the classical change $\delta\tilde{\varphi}_c$ becomes

$$\delta\tilde{\varphi}_c \simeq |\dot{\tilde{\varphi}}|H^{-1} \simeq \frac{2}{\tilde{\varphi}}. \quad (17)$$

On the other hand, the amplitude of quantum fluctuations $\delta\tilde{\varphi}_q \simeq H/(2\pi)$. Thus, the above classical equation of motion for $\tilde{\varphi}$ is valid only if $\tilde{\varphi} \ll \tilde{\varphi}_i \equiv \sqrt{4\pi\sqrt{6}/m}$. Otherwise, the universe is in a self-reproduction stage of eternal inflation [17,18] and the current horizon scale is contained in a domain where $\tilde{\varphi}$ got smaller than $\tilde{\varphi}_i$ and the classical description of $\tilde{\varphi}$ with the above classical equation of motion became feasible. Therefore, we consider only the region $\tilde{\varphi} \ll m^{-1/2}$.⁶

⁵If we take the higher order term $\xi|X|^4$ with $\xi < -9/8$ in the Kähler potential, the effective mass squared of X becomes larger than $9H^2/4$ so that X rapidly oscillates around the origin and its amplitude goes to zero.

⁶In this region we may safely neglect the higher order terms of $\Xi\Phi$ in Eq. (3).

Next, in order to estimate the amplitude of quantum fluctuations of X , we use the Fokker-Planck equation for the statistical distribution function of X , $P[X, t]$,

$$\frac{\partial}{\partial t} P[X, t] = \frac{1}{3H(t)} \frac{\partial}{\partial X} \left(m^2 X P[X, t] \right) + \frac{H^3(t)}{8\pi^2} \frac{\partial^2}{\partial X^2} P[X, t], \quad (18)$$

which is obtained through the Langevin equation based on Eq. (15) with use of the stochastic inflation method of Starobinsky [19]. Then, the time evolution of the root mean square (RMS) of fluctuations of X is given by

$$\frac{d}{dt} \langle (\Delta X)^2 \rangle = -\frac{2m^2}{3H} \langle (\Delta X)^2 \rangle + \frac{H^3}{4\pi^2}. \quad (19)$$

Taking $\tilde{\varphi}$ as a time variable in Eq. (19) by virtue of Eq. (14), we find that the RMS fluctuations of X in an initially homogeneous domain at $\tilde{\varphi} = \tilde{\varphi}_i$ are given by

$$\langle (\Delta X)^2 \rangle = \frac{m^2}{96\pi^2} \left(\tilde{\varphi}_i^2 \tilde{\varphi}^2 - \tilde{\varphi}^4 \right), \quad (20)$$

at the epoch $\tilde{\varphi}$. Taking $\tilde{\varphi}_i \simeq \sqrt{4\pi\sqrt{6}/m}$, $\langle (\Delta X)^2 \rangle$ asymptotically approaches

$$\langle (\Delta X)^2 \rangle \simeq \frac{\sqrt{6}m}{24\pi} \tilde{\varphi}^2. \quad (21)$$

On the other hand, from Eq. (16), the classical value of X , X_c is at most $\tilde{\varphi}/\tilde{\varphi}_i \simeq \sqrt{m}\tilde{\varphi}/\sqrt{4\pi\sqrt{6}}$. Thus, during the chaotic inflation,

$$\sqrt{\langle (\Delta X)^2 \rangle} \lesssim X_c \sim \sqrt{m}\tilde{\varphi} \ll 1 \ll \tilde{\varphi} = \varphi + \delta_I, \quad (22)$$

because $m \ll 1$ as shown later. Thus, for X , quantum fluctuations are smaller than the classical value, and moreover our approximation that both η and X are much smaller than unity is consistent throughout the chaotic inflation.

Let us investigate the minimum of the potential after the chaotic inflation. Since $X \sim \sqrt{m}(\varphi + \delta_I) \ll 1$ as shown above, the potential can be rewritten as⁷

$$V(\eta, \varphi, X = 0) = \frac{1}{2} m^2 \exp \left(\sqrt{2}\delta_2\eta + \eta^2 \right) \left[(\eta + \delta_R)^2 + (\varphi + \delta_I)^2 \right]. \quad (23)$$

The extreme of the potential is obtained by the conditions $\partial V/\partial\varphi = \partial V/\partial\eta = 0$,

$$\begin{aligned} \frac{\partial V}{\partial\varphi} &= m^2 \exp \left(\sqrt{2}\delta_2\eta^2 + \eta^2 \right) (\varphi + \delta_I) = 0, \\ \frac{\partial V}{\partial\eta} &= m^2 \exp \left(\sqrt{2}\delta_2\eta^2 + \eta^2 \right) \left\{ \left(\eta + \frac{\delta_2}{\sqrt{2}} \right) \left[(\eta + \delta_R)^2 + (\varphi + \delta_I)^2 \right] + (\eta + \delta_R) \right\} = 0, \end{aligned} \quad (24)$$

⁷In fact, after the inflation, the X field also decays into standard particles so that the amplitude of X rapidly goes to zero. Hence, we can safely set X to be zero.

which yields $\varphi = -\delta_I$ and

$$(\eta + \delta_R) \left\{ 2\eta^2 + (2\delta_R + \sqrt{2}\delta_2)\eta + \sqrt{2}\delta_R\delta_2 + 1 \right\} = 0. \quad (25)$$

Thus, for $|\sqrt{2}\delta_R - \delta_2| \leq 4$, $\eta = -\delta_R$ is only a minimum of the potential. Otherwise, there is another local minimum near the minimum during the inflation, which generally prevents the inflation from ending. Hence, the condition that $\delta_2 \sim |\delta_1| \lesssim m \sim 10^{-5}$ must be satisfied for a successful inflation.

Now that preparations are complete, the density fluctuations produced by this chaotic inflation is estimated as [20]

$$\frac{\delta\rho}{\rho} \simeq \frac{1}{5\sqrt{3}\pi} \frac{m}{2\sqrt{2}} \left[(\varphi + \delta_I)^2 + X^2 \right]. \quad (26)$$

Since $X \ll \varphi + \delta_I$ as shown above, the amplitude of the density fluctuations is actually determined only by the φ field. Then, the normalization at the COBE scale $[\delta\rho/\rho \simeq 2 \times 10^{-5}]$ for $(\varphi + \delta_I)_{\text{COBE}} \simeq 14$ [1] gives ⁸

$$m \simeq 10^{13} \text{ GeV} \simeq 10^{-5}. \quad (27)$$

After the inflation ends, the inflaton field φ begins to oscillate and its successive decays cause reheating of the universe. The reheating may take place by introducing the following superpotential:

$$W = \delta_4 X H_u H_d, \quad (28)$$

where $\delta_4 = g \langle \Pi \rangle$ is a constant associated with the breaking of the Z_2 symmetry. For $g = \mathcal{O}(1)$, $\delta_4 \sim |\delta_1| \lesssim m \sim 10^{-5}$ as shown above. H_u and H_d are a pair of Higgs doublets. Taking the R-charge and the Z_2 charge of $H_u H_d$ to be zero and positive, the above superpotential is invariant under the $U(1)_R$ symmetry.

Then, we have a coupling of the inflaton φ to the Higgs boson doublets as

$$L \simeq \delta_4 m \tilde{\varphi} H_u H_d, \quad (29)$$

which gives the reheating temperature T_{RH} ⁹

$$T_{\text{RH}} \lesssim 10^9 \text{ GeV} \left(\frac{\delta_4}{10^{-5}} \right) \left(\frac{m}{10^{13} \text{ GeV}} \right)^{1/2}. \quad (30)$$

Since $\delta_4 \lesssim m \sim 10^{-5}$, the reheating temperature T_{RH} becomes less than 10^9 GeV . Such a reheating temperature is low enough to avoid the gravitino problem. Recently, nonthermal

⁸The spectral index $n_s \simeq 0.96$ for $(\varphi + \delta_I)_{\text{COBE}} \simeq 14$.

⁹Field X decays into the Higgsinos \widetilde{H}_u and \widetilde{H}_d through the Yukawa interaction in Eq. (28) with the similar decay rate. Thereafter, field X rapidly goes to zero so that a pair of Higgs doublets do not acquire additional masses.

production at the preheating stage was found to be important in some inflation models [21]. For the present model, as shown by Kallosh *et al.* [21], nonthermal production of gravitinos at the preheating phase is roughly estimated as

$$\left(\frac{n_{3/2}}{s}\right)_{\text{nonTH}} \sim \frac{m^3}{m^2/T_R} \lesssim 10^{-14} \left(\frac{T_R}{10^9 \text{ GeV}}\right) \left(\frac{m}{10^{13} \text{ GeV}}\right), \quad (31)$$

where $n_{3/2}$ and s are the number density of gravitinos and entropy density. This is much less than the thermal production given by $(n_{3/2}/s)_{\text{TH}} \sim 10^{-12}(T_R/10^9 \text{ GeV})$ and hence we can neglect the nonthermal production of gravitinos.

IV. LEPTOGENESIS VIA THE INFLATON DECAY IN CHAOTIC INFLATION

In this section, we discuss the leptogenesis scenario via the inflaton decay in the above chaotic inflation model. Many leptogenesis scenarios have been proposed, so far, depending on the production mechanisms of heavy Majorana neutrinos N_i [15,22–26]. One of the most attractive scenarios is the thermal production of heavy Majorana neutrinos N_i ($i = 1 - 3$:the family index) during the reheating stage after inflation. Detailed analyses [22], however, show that enough lepton asymmetry is produced to explain the observed baryon number density only if the reheating temperature is as high as 10^{10} GeV .¹⁰ Such a high reheating temperature may cause the gravitino problem unless the gravitino mass is very light ($\lesssim 1 \text{ KeV}$) [27] or very heavy ($\gtrsim 3 \text{ TeV}$) [28].¹¹ Another interesting scenario¹² is that heavy Majorana neutrinos N_i are produced nonthermally via the decay of the inflaton [24,25]. We consider, here, a leptogenesis scenario via the inflaton decay in the above mentioned chaotic inflation model.¹³

For our purpose, we extend the Z_2 symmetry into a Z_4 symmetry. The charges of the Z_4 symmetry for various supermultiplets are given in Table I. Then, we introduce the following superpotential invariant under the $U(1)_R$ and the Z_4 symmetries:

$$W = \lambda_i m \Phi N_i N_i + \gamma_i \Pi N_i N_i, \quad (32)$$

¹⁰In our model, when $\delta_4 = g \langle \Pi \rangle \sim 10^{-4}$ with $g = \mathcal{O}(10)$, the reheating temperature T_{RH} becomes as high as 10^{10} GeV so that the thermal production of heavy Majorana neutrinos N_i leads to enough lepton asymmetry to explain the observed baryon number density.

¹¹Another solution where the gravitino is the lightest supersymmetric particle of masses from 10 to 100 GeV is proposed [29].

¹²Giudice *et al.* discussed the production of heavy Majorana neutrinos during preheating and the successive leptogenesis [24]. But, in our model, as given later in Eq. (32), the Yukawa coupling of the inflaton with heavy Majorana neutrinos is so small that sufficient lepton asymmetry cannot be produced to explain the observed baryon number density.

¹³In Ref. [30], the direct baryogenesis scenario via inflaton decay is discussed in the context of the chaotic inflationary model in $SU(1,1) \ N = 1$ supergravity [9].

where λ_i and γ_i are constants and Π is the suprior field introduced before, whose vacuum value $\langle \Pi \rangle$ leads to the breaking of the Z_4 symmetry¹⁴ and must be less than $m \sim 10^{-5}$. Here, we set $\langle \Pi \rangle \sim m \sim 10^{-5}$. The Majorana masses of right-handed neutrinos M_i is given by $M_i = \gamma_i \langle \Pi \rangle$. For $\gamma_3 = \mathcal{O}(1)$, $M_3 \sim 10^{-5} \sim 10^{13}$ GeV. The inflaton φ and the orthogonal field η can decay into right handed scalar neutrinos N_i through the above Yukawa interactions if $M_i < m/2$. Both decay rates are similar and given by

$$\Gamma_\varphi \simeq \Gamma_\eta \simeq \lambda^2 \frac{m^3}{32\pi} \sim 10\lambda^2 \text{ GeV}, \quad (33)$$

with $\lambda^2 \equiv \Sigma \lambda_i^2$ and i runs for $M_i \ll m$.¹⁵ Then, the reheating temperature T_{RH} is given by

$$T_{\text{RH}} \sim 10^9 \lambda \text{ GeV}. \quad (34)$$

For $\lambda < g$, the decay into $H_u H_d$ [see Eq. (29)] becomes the dominant decay mode of the inflaton so that the reheating temperature becomes $T_{\text{RH}} \sim 10^9 g \text{ GeV}$ and the branching ratio of the decay into right-handed neutrinos becomes $\mathcal{O}(\lambda^2/g^2)$ because $\delta_4 = g \langle \Pi \rangle \sim gm \sim 10^{-5}g$.

The produced N_i decay into leptons l_j and Higgs doublets H_u through the following Yukawa interactions of Higgs supermultiplets, which is invariant under the $U(1)_R$ and the Z_4 symmetries;

$$W = (h_\nu)_{ij} N_i l_j H_u. \quad (35)$$

Here we have taken a basis where the mass matrix for N_i is diagonal, and have assumed that quarks and leptons can be classified into the $SU(5)$ multiplets, $\mathbf{10} = (q, u^c, e^c)$, $\mathbf{5} = (d^c, l)$, and $\mathbf{1} = (N)$. We also assume $|(h_\nu)_{i3}| > |(h_\nu)_{i2}| \gg |(h_\nu)_{i1}|$ ($i = 1, 2, 3$). We consider only the decay of N_1 assuming that the mass M_1 is much smaller than the others, M_2 and M_3 . The decay of N_1 has two decay channels,

$$N_1 \rightarrow H_u + l, \quad (36)$$

$$\rightarrow \overline{H_u} + \bar{l}. \quad (37)$$

These decay channels have different branching ratios if CP symmetry is violated. Interference between the tree-level and the one-loop diagrams including vertex and self-energy corrections generates lepton asymmetry [15,31–33],

$$\begin{aligned} \epsilon_1 &\equiv \frac{\Gamma(N_1 \rightarrow H_u + l) - \Gamma(N_1 \rightarrow \overline{H_u} + \bar{l})}{\Gamma_{N_1}} \\ &= -\frac{3}{16\pi (h_\nu h_\nu^\dagger)_{11}} \left[\text{Im} (h_\nu h_\nu^\dagger)_{13}^2 \frac{M_1}{M_3} + \text{Im} (h_\nu h_\nu^\dagger)_{12}^2 \frac{M_1}{M_2} \right]. \end{aligned} \quad (38)$$

¹⁴This Z_4 symmetry is broken down to another Z_2 symmetry by $\langle \Pi \rangle \neq 0$, where this Z_2 symmetry is nothing but the so-called matter parity.

¹⁵The field X decays into \tilde{N}_i (sneutrinos) through the cross term of the superpotential with the similar decay rate to Eq. (33). φ and η also have the decay channel into \tilde{N}_i but their decay rates are much smaller than Γ_φ and Γ_η .

By use of the above hierarchy of the Yukawa coupling constants, the lepton asymmetry is dominated by the first term for $\frac{m_{\nu_3}}{m_{\nu_2}} \gtrsim \frac{M_3}{M_2}$ ¹⁶ and given by

$$\begin{aligned}
\epsilon_1 &\simeq -\frac{3\delta_{\text{eff}}}{16\pi (h_\nu h_\nu^\dagger)_{11}} \left| (h_\nu h_\nu^\dagger)_{13}^2 \right| \frac{M_1}{M_3} \\
&\simeq -\frac{3\delta_{\text{eff}}}{16\pi} |(h_\nu)_{33}|^2 \frac{M_1}{M_3} \\
&\simeq -\frac{3\delta_{\text{eff}}}{16\pi} \frac{m_{\nu_3} M_1}{\langle H_u \rangle^2} \\
&\sim -10^{-5} \delta_{\text{eff}} \left(\frac{M_1}{10^{11} \text{ GeV}} \right), \tag{39}
\end{aligned}$$

where δ_{eff} is a parameter representing the magnitude of the CP violation, m_{ν_3} is estimated by the seesaw mechanism [14] as

$$\begin{aligned}
m_{\nu_3} &\simeq \frac{|(h_\nu)_{33}|^2 \langle H_u \rangle^2}{M_3} \\
&\sim 10^{-2} \text{ eV} \left(\frac{|(h_\nu)_{33}|}{10^{-1}} \right) \left(\frac{10^{13} \text{ GeV}}{M_3} \right), \tag{40}
\end{aligned}$$

which is consistent with the mass suggested from the Super-kamiokande experiments [13] for $|(h_\nu)_{33}| \sim 10^{-1}$ and $M_3 \sim 10^{13} \text{ GeV}$.

The total decay rate of N_1 , Γ_{N_1} , is given by

$$\begin{aligned}
\Gamma_{N_1} &= \Gamma(N_1 \rightarrow H_u + l) + \Gamma(N_1 \rightarrow \overline{H_u} + \bar{l}) \\
&\simeq \frac{1}{8\pi} \sum |(h_\nu)_{1i}|^2 M_1 \\
&\simeq \frac{1}{8\pi} |(h_\nu)_{13}|^2 M_1 \\
&\sim 10^5 \text{ GeV} \left(\frac{|(h_\nu)_{13}|}{10^{-2}} \right)^2 \left(\frac{M_1}{10^{11} \text{ GeV}} \right). \tag{41}
\end{aligned}$$

Thus, for a wide range of parameters, the decay rate Γ_{N_1} is much larger than the decay rate of the inflaton Γ_φ so that the produced N_1 immediately decays into leptons and Higgs supermultiplets.

Before estimating the lepton asymmetry produced in our model, let us evaluate the lepton asymmetry needed to explain the observed baryon number density. A part of produced lepton asymmetry, exactly speaking, $B - L$ asymmetry is converted into baryon asymmetry through the sphaleron processes, which can be estimated as [34]

$$\frac{n_B}{s} \simeq -\frac{8}{23} \frac{n_L}{s}, \tag{42}$$

¹⁶Even if the second term dominates, the discussion also runs parallel.

where we have assumed the standard model with two Higgs doublets and three generations. In order to explain the observed baryon number density,

$$\frac{n_B}{s} \simeq (0.1 - 1) \times 10^{-10}, \quad (43)$$

we need the lepton asymmetry,

$$\frac{n_L}{s} \simeq -(0.3 - 3) \times 10^{-10}. \quad (44)$$

Now we estimate the lepton asymmetry produced through the inflaton decay. For $M_1 \gtrsim 10^{11} \lambda$ GeV, M_1 is one hundred times larger than the reheating temperature T_{RH} . In this case, the produced N_1 is never in thermal equilibrium. Then, the ratio of the lepton number to entropy density can be estimated as

$$\begin{aligned} \frac{n_L}{s} &\simeq \frac{3}{2} \epsilon_1 B_r \frac{T_R}{m} \\ &\sim -10^{-7} \delta_{\text{eff}} B_r \left(\frac{T_R}{10^9 \text{ GeV}} \right) \left(\frac{M_1}{m} \right) \\ &\sim -10^{-9} \delta_{\text{eff}} B_r \left(\frac{T_R}{10^9 \text{ GeV}} \right) \left(\frac{M_1}{10^{11} \text{ GeV}} \right) \left(\frac{10^{13} \text{ GeV}}{m} \right), \end{aligned} \quad (45)$$

where B_r is the branching ratio of the inflaton decay into N_1 . For $M_3 \sim M_2 \sim m \sim 10^{13}$ GeV, the decay into N_3 and N_2 are prohibited kinematically or suppressed by the phase space and hence $B_r = \mathcal{O}(1)$ for $\lambda_1 = \mathcal{O}(1)$. In this case, we obtain $T_{\text{RH}} \sim 10^9$ GeV, which results in $n_L/s \sim -10^{-9} \delta_{\text{eff}}$. Thus, our model of leptogenesis works well for $\gamma_2 \simeq \gamma_3 = \mathcal{O}(1)$, $\delta_{\text{eff}} = \mathcal{O}(1)$, and $\lambda_1 = \mathcal{O}(1)$ [see Eq. (32)].

Finally, we make a comment on the Froggatt-Nielsen (FN) mechanism [35] based on a spontaneously broken $U(1)_F$ family symmetry, which gives a natural explanation for the observed mass hierarchy in mass matrices of quarks and charged leptons. The $U(1)_F$ symmetry is broken by a gauge singlet scalar field Δ with FN charge $Q_\Delta = -1$, whose condensation $\langle \Delta \rangle$ gives rise to the Yukawa coupling constants. That is, the Yukawa couplings of Higgs supermultiplets are given through nonrenormalizable interactions with Δ ,

$$W = g_{ij} \Delta^{Q_i + Q_j} \Psi_i \Psi_j H_{u(d)}, \quad (46)$$

where Q_i are the FN charges of quark and lepton supermultiplets Ψ_i , g_{ij} are coupling constants of the order unity, and H_u, H_d are Higgs supermultiplets with FN charges zero. In particular, $(h_\nu)_{ij} = g_{ij} \langle \Delta \rangle^{Q_{N_i} + Q_{l_j}}$. Then, the observed mass hierarchy can be well explained if we take $\epsilon \equiv \langle \Delta \rangle \simeq 1/17$ and the FN charges of quark and lepton supermultiplets shown in Table II [36].

If the Froggatt-Nielsen mechanism is adopted, the above discussion on the leptogenesis also holds except for three points. First of all, two contributions to the lepton asymmetry in Eq. (38) become comparable. Next, the coupling constants γ_i in Eq. (32) becomes $\gamma_3 \sim \gamma_2 = \mathcal{O}(1)$ and $\gamma_1 = \mathcal{O}(10^{-2})$. Therefore, M_3 and $M_2 \sim 10^{13}$ GeV automatically become comparable with the mass of inflaton φ and other fields η and X , i.e. $\sim m \sim 10^{13}$ GeV, so that the decay into N_3 and N_2 are prohibited kinematically or suppressed by the phase

space. Finally, λ_1 in Eq. (32) becomes $\lambda_1 = \mathcal{O}(10^{-2})$ so that the reheating temperature T_{RH} becomes 10^7 GeV. In this case, unless $g < \mathcal{O}(10^{-2})$, the decay mode into the Higgs doublet in Eq. (29) must be forbidden because otherwise the branching ratio becomes small as $B_r \sim \lambda_1^2/g^2$ and the produced lepton asymmetry may be too small. If, for example, we set the R-charge of $H_u H_d$ to be nonzero,¹⁷ the superpotential in Eq. (28) is prohibited. Then, the ratio of lepton number density to entropy density can be estimated as

$$\frac{n_L}{s} \sim -10^{-11} \left(\frac{T_R}{10^7 \text{ GeV}} \right) \left(\frac{M_1}{10^{11} \text{ GeV}} \right) \left(\frac{10^{13} \text{ GeV}}{m} \right), \quad (47)$$

which is marginally consistent with the baryon number density in the present universe.

V. DISCUSSION AND CONCLUSIONS

In this paper we have comprehensively investigated a natural chaotic inflation model with the shift symmetry in supergravity. In particular, the forms of the Kähler potential and the superpotential have been discussed. In order to suppress higher order terms of the inflaton field in the superpotential, the shift symmetry is extended into that including the suprior field Ξ with the combination $\Xi\Phi$ invariant. Also, the linear term of X in the superpotential is suppressed by introducing the Z_2 symmetry. We have found that if the magnitude of the breaking of the Z_2 symmetry is equal or smaller than that of the shift symmetry, a desired chaotic inflation can take place.

We have also discussed the leptogenesis via the inflaton decay in this chaotic inflation model. The inflaton φ can decay into right-handed neutrinos through the Yukawa interactions suppressed by the breaking of the shift symmetry, which leads to low reheating temperature enough to avoid the overproduction of gravitinos. Right-handed neutrinos acquire their masses associated with the breaking of a Z_4 symmetry which is an extension of the Z_2 symmetry, whose magnitude is consistent with the result from the Super-kamiokande experiment. Then, we have found that for a wide range of parameters, the lepton asymmetry enough to explain the observed baryon number density is produced. Also, when the Froggatt-Nielsen mechanism is adopted as the mechanism to explain the hierarchy for the masses of leptons and quarks, we have obtained the lepton asymmetry, which is marginally consistent with the baryon number density in the present universe.

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¹⁷In this case, other necessary Yukawa interactions are all permitted, taking the R-charges of H_u , H_d , $\mathbf{5}^*$, and $\mathbf{10}$ to be $2a/5$, $3a/5$, $1 - 2a/5$, and $1 - a/5$, where the R-charge of $H_u H_d$ is a .

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TABLES

	Φ	X	Ξ	Π	N	H_u	H_d	$\mathbf{5}^*$	$\mathbf{10}$
Q_R	0	2	0	0	1	0	0	1	1
Z_4	2	2	0	2	1	2	2	1	1

TABLE I. The charges of various supermultiplets of $U(1)_R \times Z_4$. Here, R charge of $H_u H_d$ is assigned to be 0. All supermultiplets of quarks and leptons have the Z_4 charge 1 and Higgs supermultiplets H_u and H_d carry the Z_4 charge 2.

Ψ_i	$\mathbf{5} = (d^c, l)$			$\mathbf{10} = (q, u^c, e^c)$			$\mathbf{1} = (N)$		
	$\mathbf{5}_3$	$\mathbf{5}_2$	$\mathbf{5}_1$	$\mathbf{10}_3$	$\mathbf{10}_2$	$\mathbf{10}_1$	$\mathbf{1}_3$	$\mathbf{1}_2$	$\mathbf{1}_1$
Q_i	1	1	2	0	1	2	0	0	1

TABLE II. The FN charges of quark and lepton supermultiplets assumed throughout this paper.